

# The moduli and gravitino (non)-problems in models with strongly stabilized moduli

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**Abstract.** In gravity mediated models and in particular in models with strongly stabilized moduli, there is a natural hierarchy between gaugino masses, the gravitino mass and moduli masses:  $m_{1/2} \ll m_{3/2} \ll m_\phi$ . Given this hierarchy, we show that 1) moduli problems associated with excess entropy production from moduli decay and 2) problems associated with moduli/gravitino decays to neutralinos are non-existent. Placed in an inflationary context, we show that the amplitude of moduli oscillations are severely limited by strong stabilization. Moduli oscillations may then never come to dominate the energy density of the Universe. As a consequence, moduli decay to gravitinos and their subsequent decay to neutralinos need not overpopulate the cold dark matter density.

**Keywords:** dark matter theory, supersymmetry and cosmology, particle physics - cosmology connection

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## 1 Introduction

The presence of light weakly-interacting fields (moduli) in the early universe has problematic consequences in cosmology [1–4]. Their late decay implies that these moduli may eventually dominate the energy density of the universe, redshifting as nonrelativistic pressureless matter. When they decay, they generally generate too much entropy, diluting any baryon asymmetry generated at earlier times, while failing to reheat the universe sufficiently to restart nucleosynthesis. Furthermore, their late decays may lead to the overproduction of the lightest supersymmetric particle (LSP) [5, 6].

A well known example of this problem appears in the context of the Polonyi model of soft supersymmetry breaking in  $\mathcal{N} = 1$  supergravity [7]. At the supersymmetry breaking minimum, the scalar Polonyi field  $Z$  has a mass of the order of the gravitino mass,  $m_Z \sim m_{3/2}$ . Since it couples with gravitational strength to matter fields, its decay rate is  $\Gamma_Z \sim m_{3/2}^3/M_P^2$ , where  $M_P \simeq 2.4 \times 10^{18}$  GeV denotes the reduced Planck mass. Moreover, during an inflationary epoch, the Polonyi field will be displaced from the minimum of the potential. Generically, this displacement is of the order of the Planck scale,  $\Delta Z \sim \mathcal{O}(M_P)$  [8]. The combination of a large initial displacement, a small mass and a small decay rate is at the root of the Polonyi problem. The resulting reheating temperature,

$$T_R(Z) \sim \frac{m_{3/2}^{3/2}}{M_P^{1/2}}, \quad (1.1)$$

is smaller than the temperature required by nucleosynthesis  $T_N \sim 1$  MeV unless  $m_{3/2} \gtrsim 10$  TeV. Even more problematic is the entropy release from the decay of Polonyi oscillations [1]

$$\frac{s_f}{s_i} \sim \frac{M_P}{m_{3/2}}, \quad (1.2)$$

where  $s_{i,f}$  are the entropy densities before and after decay. This late injection of entropy would severely dilute any pre-existing baryon asymmetry.

When the Polonyi field decays into lighter supersymmetric particles, the eventual overproduction of the LSP is likely. In particular, gravitinos may be copiously produced by the decay of the Polonyi field leading to a gravitino problem [9, 10]. If the gravitino is unstable, the decay products will in turn eventually decay into the LSP, generically resulting in a dark matter relic density much larger than that observed [5, 6]. Furthermore, the gravitino decay rate is also of the order  $\Gamma_{3/2} \sim m_{3/2}^3/M_P^2$ . So it too can decay late and cause problems for nucleosynthesis.

Some of the problems of the Polonyi field can be solved by giving it a larger mass [4, 5, 11]. For a modulus mass  $m_Z$  (or  $m_{3/2}$ ) larger than  $\mathcal{O}(10 \text{ TeV})$ , the reheating temperature is high enough to restart nucleosynthesis after decay. The late time entropy release would nevertheless dilute the results of any previous nucleosynthesis or baryogenesis. If the baryon asymmetry is generated by a very effective mechanism, such as the Affleck-Dine (AD) mechanism [12], which can generate a baryon-to-entropy ratio as large as  $\mathcal{O}(1)$ , the resulting increase in entropy could provide the necessary dilution factor to yield the observed baryon asymmetry [13, 14]. Other potential solutions to the Polonyi problem have also been discussed [15–21].

In this paper, we will consider a strongly stabilized hidden sector with a Polonyi type superpotential as the source of soft supersymmetry breaking. This non-minimal Polonyi model was first introduced in [22, 23], and later used as part of the so-called O’KKLT mechanism [24–26]. There are also several recent phenomenological studies of strongly stabilized moduli [27–35]. In [36], the evolution of the strongly stabilized Polonyi sector was studied in the context of a realization of chaotic inflation and the AD mechanism. In all of these models the cosmological problems are addressed by generating a hierarchy between the Polonyi and gravitino masses,  $m_Z \gg m_{3/2}$ , effectively stabilizing the Polonyi field during inflation. Furthermore, the gravitino mass may be made hierarchically larger than the weak scale as in models [27–33]. Although this hierarchy between the gravitino mass and the weak scale is useful in curing some cosmological problems, it is also motivated by the large Higgs mass seen at the LHC [37, 38]. This hierarchy also implies a decay rate for the gravitino which is much larger than in the standard Polonyi model.

For the mass scale associated with strong stabilization, we will assume  $\Lambda \ll M_P$ . This new mass scale not only fixes the hierarchy between the Polonyi field and the gravitino, but also fixes the vacuum expectation value (vev) of the Polonyi field to a value much smaller than the Planck scale, thus reducing the amplitude of oscillations and hence the energy stored in the Polonyi field. Here, we will derive an upper limit on  $\Lambda$  by requiring that the Polonyi field never comes to dominate the energy density of the Universe. For a sufficiently large mass for the Polonyi field *and* restricted vev, dilution of the products of nucleosynthesis and baryogenesis can be avoided. In addition, as we will show, the relic density of supersymmetric cold dark matter resulting from the decay of the modulus can be made consistent with current observations [39]. Unlike the scenarios in which the energy density is dominated by the modulus [40, 41], the production of lightest supersymmetric particles does not proceed directly, but through the intermediate decay to gravitinos. By further requiring that the LSP does not over-close the Universe, we are able to derive yet a stronger bound on the stabilization scale,  $\Lambda \lesssim 10^{-3} M_P$ .

Our paper is organized as follows: in the next section, we describe the mechanism for strong stabilization and discuss the dominant possible decay modes of the inflaton. In section 3, we incorporate an inflationary background to describe the evolution of the scalar fields. Here we will derive the conditions on strong stabilization such that the Polonyi field or moduli never come to dominate the energy density of the universe and hence lead to an insignificant increase in the entropy density. In section 4, we determine the resulting non-thermal relic density of LSPs and show the conditions under which it is compatible with observations. Our summary and conclusions are given in section 5.

## 2 Strong stabilization and decay modes

In what follows, we will distinguish between three sectors of the theory: the supersymmetry breaking sector characterized by a Polonyi-like field  $Z$ ; an inflation sector characterized by an

inflaton,  $\eta$ ; and a matter sector characterized generically by a set of fields,  $\phi$ . The strongly stabilized Polonyi sector is described by a superpotential

$$W = \mu^2(Z + \nu), \quad (2.1)$$

where the parameter  $\nu$  is adjusted so that the cosmological constant vanishes at the supersymmetry breaking minimum. Unless explicitly noted, we will work in units where the reduced Planck mass,  $M_P$ , has been set to be unity. The Kähler potential includes a strongly stabilizing term added to the minimal term [22, 23],

$$K = Z\bar{Z} - \frac{(Z\bar{Z})^2}{\Lambda^2}, \quad (2.2)$$

where it is assumed that the mass scale  $\Lambda \ll 1$ . For simplicity, we will denote by  $Z$  both the chiral superfield and its scalar component. The scalar potential derived from the Kähler potential (2.2) and the superpotential (2.1) is given by [42–46]

$$\begin{aligned} V &= e^K (K^{Z\bar{Z}} D_Z W \bar{D}_{\bar{Z}} \bar{W} - 3|W|^2) \\ &= \mu^4 e^{Z\bar{Z} - (Z\bar{Z})^2/\Lambda^2} \left[ \frac{|1 + \bar{Z}(1 - 2Z\bar{Z}/\Lambda^2)(Z + \nu)|^2}{1 - 4(Z\bar{Z})/\Lambda^2} - 3|Z + \nu|^2 \right], \end{aligned} \quad (2.3)$$

where

$$D_Z W = (\partial_Z K)W + \partial_Z W. \quad (2.4)$$

In order to obtain phenomenologically acceptable soft scalar masses, we assume that this Polonyi sector is hidden from the visible sector. Therefore, if  $\phi$  denotes collectively the matter superfields, the Kähler potential and the superpotential are assumed to be separable

$$K = K(Z, \bar{Z}) + K(\phi, \bar{\phi}), \quad (2.5)$$

$$W = W(Z) + W(\phi). \quad (2.6)$$

$K(Z, \bar{Z})$  is assumed to be given by (2.2) and for our purposes here, it is sufficient to assume that  $K(\phi, \bar{\phi}) = \phi\bar{\phi}$ . We note that phenomenological studies [31–33] show that the visible sector often requires a Giudice-Masiero term of the form

$$K(\phi, \bar{\phi}) \supset c_H H_1 H_2 + \text{h.c.}, \quad (2.7)$$

where  $H_1, H_2$  are the Higgs fields of the MSSM and  $c_H$  is some dimensionless constant.

Strongly stabilized models of the type we are considering here tend to have small supersymmetric breaking  $A$ -terms. Their tree level values are highly suppressed with  $A \propto m_{3/2}\Lambda^2/M_P^2$  [31] and as a result the Higgs mass determined at the LHC tends to require rather large sfermion masses. Furthermore, if the gauge kinetic function is independent of  $Z$ , the tree-level gaugino masses are vanishing and the leading order contribution is the one-loop anomaly mediated contribution. These types of models have been considered on phenomenological grounds with tree-level sfermion masses and loop suppressed gaugino masses [31–33]. Although this is the spectrum we have implicitly assumed, it is possible that the gauginos get tree-level masses and that  $A$ -terms are generated from some additional interaction in the Kähler potential. In this case, the LHC constraints on sfermion masses would be weaker. We will often normalize the constraints on  $\Lambda$  in terms of  $m_{3/2} \sim 100 \text{ TeV}$ , since this is the

sfermion mass spectrum we have implicitly assumed. If, however, we take a lighter gravitino mass, the constraints on  $\Lambda$  will be strengthened, but still with an acceptable order of magnitude.

The complex field  $Z$  can be parametrized in terms of its real and imaginary parts,

$$Z = \frac{1}{\sqrt{2}}(z + i\chi). \quad (2.8)$$

In this parametrization, the supersymmetry breaking Minkowski minimum is found to be real and located at

$$\langle z \rangle_{\text{Min}} \simeq \frac{\Lambda^2}{\sqrt{6}}, \quad \langle \chi \rangle = 0, \quad \nu \simeq \frac{1}{\sqrt{3}}, \quad (2.9)$$

for  $\Lambda \ll 1$ . The supersymmetry breaking mass scale given by the gravitino mass is

$$m_{3/2} = \langle e^{K/2} W \rangle \simeq \mu^2 / \sqrt{3}, \quad (2.10)$$

whereas the mass squared of both  $z$  and  $\chi$  are

$$m_{z,\chi}^2 \simeq \frac{12 m_{3/2}^2}{\Lambda^2} \gg m_{3/2}^2. \quad (2.11)$$

Thus, for  $\Lambda \ll 1$ , we obtain the hierarchy mentioned earlier between the modulus and the gravitino. Note that for consistency, we must not take  $\Lambda$  so small so that  $m_Z > \Lambda$ . This imposes

$$\Lambda > 6 \times 10^{-7} M_P \left( \frac{m_{3/2}}{10^{-13} M_P} \right)^{1/2}. \quad (2.12)$$

The goldstino is the fermionic component of  $Z$ ,  $\psi_Z$ . In the unitary gauge it is absorbed by the gravitino, becoming its longitudinal component, via the super-Higgs mechanism [47]. It is worth noting that we have also explored the scenario in which the non-minimal Kähler term is positive,  $K = Z\bar{Z} + (Z\bar{Z})^2/\Lambda^2$ . In this case, for  $\Lambda \ll 1$ , along the real axis ( $\chi = 0$ ), the parameter  $\nu$  can be tuned to yield a Minkowski supersymmetry breaking minimum, together with an Anti-de Sitter minimum, both separated by a barrier of finite size about the origin. However, as a function of real and imaginary parts  $z, \chi$ , the AdS extremum is found to be the global minimum, while the Minkowski extremum is actually a saddle point, and is connected to the global minimum in the complex direction.

The decay modes of the Polonyi field are determined by its couplings to matter and gauge fields, and to the gravitino. The interaction with matter scalars follows from the Lagrangian

$$\mathcal{L}_S = G_{i\bar{j}} D_\mu \phi_i D^\mu \bar{\phi}_{\bar{j}} - e^G (G_i G^{i\bar{j}} G_{\bar{j}} - 3), \quad (2.13)$$

with  $G = K + \log |W|^2$  the Kähler function. Eq. (2.5) implies that the relevant interaction for decay corresponds to the potential term. As noted earlier, for simplicity, we consider a minimal Kähler potential for matter fields. Under the assumption that the vevs of the matter fields are either zero or at most of order the weak scale, we set  $K_i, W_i \ll 1$ . The scalar potential can then be Taylor expanded to give the two body decay coupling,

$$\mathcal{L}_{S,2} = \sqrt{3} m_{3/2} (m_{3/2} - \bar{W}(\bar{\phi})) Z \phi_i \bar{\phi}_{\bar{i}} + \text{h.c.} + \mathcal{O}(\Lambda^2). \quad (2.14)$$

The resulting decay rate is suppressed by  $\Lambda \ll 1$ . Restoring the Planck mass  $M_P$ , the width is

$$\Gamma(z \rightarrow \phi_i \bar{\phi}_{\bar{i}}) \simeq \frac{\sqrt{3} \Lambda}{32\pi} \frac{m_{3/2}^3}{M_P^3}. \quad (2.15)$$

A further expansion reveals that the three body decays to matter scalars are determined by the Yukawa couplings of the matter fields,

$$\mathcal{L}_{S,3} = \sqrt{3}m_{3/2}W_{ijk}\bar{Z}\phi_i\phi_j\phi_k + \text{h.c.} + \mathcal{O}(\Lambda^2). \quad (2.16)$$

In this case the decay rate is enhanced by  $\Lambda$ ,

$$\Gamma(Z \rightarrow \bar{\phi}_i\bar{\phi}_j\bar{\phi}_k) \simeq \frac{3\sqrt{3}|W_{ijk}|^2 m_{3/2}^3}{256\pi^3\Lambda M_P}. \quad (2.17)$$

The interaction of  $Z$  with matter fermions is determined by the kinetic and mass terms of the supergravity Lagrangian,

$$\begin{aligned} \mathcal{L}_F = & \frac{i}{2}G_{ij}\bar{\chi}_R^i\gamma^\mu D_\mu\chi_R^j + \frac{i}{2}\left(-G_{ij\bar{k}} + \frac{1}{2}G_{i\bar{k}}G_j\right)\bar{\chi}_R^i\gamma^\mu D_\mu\phi^j\chi_R^k \\ & - \frac{1}{2}e^{G/2}(-G_{ij} - G_iG_j + G_{ij\bar{k}}G^{\bar{k}l}G_l)\bar{\chi}_R^i\chi_L^j + \text{h.c.} \end{aligned} \quad (2.18)$$

After the Goldstino component is subtracted out, the interactions for the two body decays  $Z \rightarrow \bar{\chi}^i\chi^j$  are found to be given by

$$\mathcal{L}_{F,2} = i\frac{\sqrt{3}}{4}\bar{\chi}_R^i\gamma^\mu\partial_\mu Z\chi_{iR} + \frac{1}{\sqrt{3}m_{3/2}}\bar{\chi}_R^i(ZW_iW_j)\chi_L^j + \frac{\sqrt{3}}{4}m_{3/2}c_H Z\tilde{H}_L\tilde{H}_R + \text{h.c.} + \mathcal{O}(\Lambda^2) \quad (2.19)$$

where  $\tilde{H}^T = [\tilde{H}_1^T, \tilde{H}_2^\dagger]$ . The squared amplitude for first term is suppressed by the masses of the final-state fermions, in addition to a factor of  $\mathcal{O}(\Lambda)$ . The amplitude of the second term is suppressed by the expectation values  $\langle W_i \rangle \ll 1$ . The interference term between the first two terms vanishes. The third term, turns out to be the most dominant two-body decay mode to matter fields. It gives a decay width of

$$\Gamma(Z \rightarrow \tilde{H}_1\tilde{H}_2) = \frac{3c_H^2}{256\pi}\left(\frac{m_{3/2}}{M_P}\right)^2 m_Z = \frac{3c_H^2}{256\pi}\frac{m_{3/2}^3}{M_P^2}\frac{M_P}{\Lambda}. \quad (2.20)$$

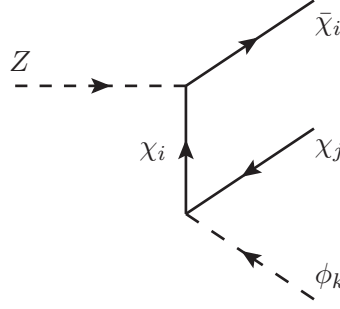
Three body decays which include fermions in the final state, and which proceed through four point vertices are also suppressed by the expectation values of  $W_i$ . The largest contribution to the decay into fermions, other than possibly the Higgsinos, is given by the fermion exchange diagram of figure 1, with the rate

$$\Gamma(Z \rightarrow \bar{\chi}_i\chi_j\phi_k) \simeq \frac{9\sqrt{3}|W_{ijk}|^2}{2048\pi^3\Lambda^3}\log\left(\frac{m_z}{m_k}\right)m_{3/2}^3M_P. \quad (2.21)$$

Here we have neglected the masses of the final state fermions.

Without an explicit coupling through the gauge kinetic function,  $Z$  does not decay into gauge bosons or gauginos at tree level. Nevertheless, the Polonyi field can still decay into the gauge supermultiplets through anomaly mediated effects [48, 49]. The corresponding decay rate is suppressed by a factor of  $\mathcal{O}(\Lambda)$ ,

$$\Gamma(Z \rightarrow gg, \tilde{g}\tilde{g}) \sim \frac{N_g\alpha^2}{256\pi^3}|K_Z|^2\frac{m_z^3}{M_P^2} \sim \frac{\Lambda N_g\alpha^2}{256\pi^3}\frac{m_{3/2}^3}{M_P^3}. \quad (2.22)$$



**Figure 1.** Three body decay of  $Z$  into a final state with two fermions

In addition to the decays into matter and gauge fields, the Polonyi field can decay into gravitinos. This process is mediated by the interaction terms

$$\begin{aligned}\mathcal{L}_{3/2} &= \frac{1}{8}\epsilon^{\mu\nu\rho\sigma}\bar{\psi}_\mu\gamma_\nu\psi_\rho G_i D_\sigma\phi_i + \frac{i}{2}e^{G/2}\bar{\psi}_{\mu L}\sigma^{\mu\nu}\psi_{\nu R} + \text{h.c.} \\ &= \frac{\sqrt{3}}{8}\epsilon^{\mu\nu\rho\sigma}\bar{\psi}_\mu\gamma_\nu\psi_\rho\partial_\sigma Z - \frac{\sqrt{3}}{8}m_{3/2}Z\bar{\psi}_\mu[\gamma^\mu, \gamma^\nu]\psi_\nu + \text{h.c.} + \mathcal{O}(\Lambda^2) + \dots\end{aligned}\quad (2.23)$$

where the ellipsis includes higher order terms in  $Z$ , as well as the couplings of the gravitino to scalar fields other than  $Z$ . The rate obtained from these couplings is enhanced by a factor of  $\Lambda^{-5}$ ,

$$\Gamma(Z \rightarrow \psi_{3/2}\psi_{3/2}) \simeq \frac{3\sqrt{3}m_{3/2}^3 M_P^3}{\pi\Lambda^5}. \quad (2.24)$$

Alternatively, the decay rate may be computed in the goldstino picture [35, 50, 51] with the same result. This rate differs from the rate computed in the standard Polonyi scenario without stabilization [48, 49], since the interaction for the goldstino comes from the strongly stabilizing contribution to the Kähler potential

$$\int d^4\theta K \supset - \int d^4\theta \frac{|Z|^4}{\Lambda^2} \supset -2 \frac{F_Z^\dagger Z^\dagger \psi_z \psi_z}{\Lambda^2} = -2\sqrt{3} \frac{m_{3/2} M_P}{\Lambda^2} Z^\dagger \psi_z \psi_z. \quad (2.25)$$

Using this and calculating the decay width we get exactly what the result above for the decays to gravitinos. The dominant channel for the spontaneous decay of the Polonyi field is therefore that to gravitinos,  $\Gamma_z^{(\text{total})} \simeq \Gamma(Z \rightarrow \psi_{3/2}\psi_{3/2})$ .

### 3 Post-inflationary dynamics

During an inflationary epoch, the scalar field  $Z$  will be displaced from its true minimum given by eq. (2.9) to smaller values. In supergravity, large masses of the order of the Hubble parameter during inflation,  $H_I$ , are generically induced on scalar fields, due to the exponential factor in (2.3). If  $\eta$  denotes the scalar field responsible for inflation, a contribution

$$\Delta V(Z) \sim e^{K(Z)} V(\eta) = c H_I^2 Z \bar{Z} + \dots \quad (3.1)$$

will typically arise, where  $c \sim 3$  in general. If this is the case, the expectation value of  $Z$  will be several orders of magnitude smaller than the true minimum,  $\langle Z \rangle_{\text{inf}} \ll \langle Z \rangle_{\text{Min}}$ . In

particular, the addition of a contribution of the form of eq. (3.1) to the potential (2.3) results in the vacuum expectation value during inflation [36]

$$\langle z \rangle_{\text{inf}} \simeq \frac{\Lambda^2}{\sqrt{6}} \left( 1 + \frac{3c}{2} \frac{H_I^2 \Lambda^2}{\mu^4} \right)^{-1} \simeq \sqrt{\frac{2}{3}} \frac{\mu^4}{3cH_I^2} \ll \langle z \rangle_{\text{Min}}. \quad (3.2)$$

At the end of inflation, the universe is dominated by the oscillations of the inflaton which leads to an expansion rate characterized by matter domination. During this period, the Hubble parameter decreases and the Polonyi field will adiabatically track the instantaneous minimum [20, 21] until the Hubble parameter becomes of the order of the mass of  $z$ ; more precisely when  $H = \frac{2}{3}m_z$ . When this occurs,  $z$  will start damped oscillations about the true supersymmetry breaking minimum (2.9). This may occur either before or after the inflaton oscillations have decayed. Thus the amplitude of oscillations in this strongly stabilized model is reduced relative to the standard case by the fact that the final vev is of order  $\Lambda^2/M_P \ll M_P$ .

The energy density and the Hubble parameter during the epoch where inflaton oscillations dominated universe can be written as

$$\rho_\eta = \frac{4}{3} m_\eta^2 M_P^2 \left( \frac{R_\eta}{R} \right)^3, \quad (3.3)$$

$$H = \frac{2}{3} m_\eta \left( \frac{R_\eta}{R} \right)^{3/2}, \quad (3.4)$$

where  $R_\eta$  denotes the cosmological scale factor at the onset of oscillations of  $\eta$ . Therefore, the oscillation of  $z$  starts when the scale factor is

$$R_z \simeq \left( \frac{\Lambda}{2\sqrt{3}} \frac{m_\eta}{m_{3/2} M_P} \right)^{2/3} R_\eta, \quad (3.5)$$

where we have used (2.11) for  $m_z$ .

During inflation, the imaginary part of  $Z$ ,  $\chi$ , at the minimum is not displaced and so evolves to the minimum of the potential and does not oscillate thereafter. Hence, the energy density of the Polonyi field is stored in the oscillations of the real part  $z$ ,

$$\rho_z \simeq \frac{1}{2} m_z^2 \langle z \rangle_{\text{Min}}^2 \left( \frac{R_z}{R} \right)^3. \quad (3.6)$$

The amplitude of the oscillations is therefore suppressed by  $\Lambda^2$ . In addition, note that for the strongly stabilized modulus we also have an enhanced mass and therefore an enhanced decay rate. Therefore, as we show below, for a sufficiently small  $\Lambda$ , the cosmological problems for  $Z$  are averted. The details of the evolution of  $Z$  depend on whether the Polonyi field decays before or after reheating. Some numerical results for the evolution of  $z$  and  $\chi$  can be found in [36].

Assuming that the inflaton decays due to gravitational-strength interactions, we can parametrize the coupling by  $d_\eta$ , such that the decay rate is

$$\Gamma_\eta = d_\eta^2 \frac{m_\eta^3}{M_P^2}. \quad (3.7)$$



In the instantaneous approximation, the inflaton decays when  $\Gamma_\eta = \frac{3}{2}H$ , or  $R_{d\eta}/R_\eta = (M_P/d_\eta m_\eta)^{4/3}$ . Comparing  $R_{d\eta}$  with  $R_z$ , we see that Polonyi oscillations begin before inflaton decay so long as

$$\left(\frac{d_\eta^2 \Lambda}{m_{3/2}}\right)^{2/3} \left(\frac{m_\eta}{M_P}\right)^2 < 1. \quad (3.8)$$

For  $m_\eta \sim 10^{-5}M_P$  and  $m_{3/2} \sim 10^{-13}M_P$ , this condition is valid for  $d_\eta^2 \Lambda < 10^2 M_P$  as we will assume. If  $\Lambda$  is very small,  $Z$  will decay before the inflaton. Very early decays of the Polonyi field will occur when  $R_{dz} < R_{d\eta}$  where  $R_{dz}$  is the scale factor at the time of  $z$  decay. As we will see, this condition is satisfied when  $d_\eta^{2/5} \Lambda/M_P < 10^{-5}$ . We will return to this case below.

After inflaton decay, the universe is filled with the relativistic decay products, with energy density and Hubble parameter

$$\rho_r = \frac{4}{3} d_\eta^{-4/3} m_\eta^{2/3} M_P^{10/3} \left(\frac{R_\eta}{R}\right)^4, \quad (3.9)$$

$$H_r = \frac{2}{3} d_\eta^{-2/3} m_\eta^{1/3} M_P^{2/3} \left(\frac{R_\eta}{R}\right)^2. \quad (3.10)$$

The corresponding reheating temperature is

$$T_R = d_\eta \left(\frac{40}{\pi^2 g_\eta}\right)^{1/4} \frac{m_\eta^{3/2}}{M_P^{1/2}}, \quad (3.11)$$

where  $g_\eta = g(T_R)$  is the effective number of degrees of freedom at reheating. For large  $\Lambda$  (but still  $\lesssim M_P$ ), the Universe may become dominated by  $z$  oscillations before they decay (as in the standard Polonyi scenario). If this should happen, the Hubble parameter becomes

$$H_z = \frac{1}{6} m_z \frac{\Lambda^2}{M_P^2} \left(\frac{R_z}{R}\right)^{3/2}. \quad (3.12)$$

In this case, the scale factor at  $z$  decay is

$$R_{dz} = \left(\frac{\pi}{6}\right)^{2/3} \frac{\Lambda^4}{m_{3/2}^{4/3} M_P^{8/3}} R_z. \quad (3.13)$$

Using (3.9) for the energy density in radiation (subdominant), (3.6) for the energy density in  $z$  oscillations, and (3.5) to relate the scale factors  $R_z$  and  $R_\eta$  we can compute the entropy increase due to  $z$  decays. Taking the entropy density in radiation to be  $s_r = 4/3(g_\eta \pi^2/30)^{1/4} \rho_r^{3/4}$ , and a similar expression for the entropy density produced from  $z$  decays, we find,

$$\frac{s_z}{s_r} \simeq 0.05 d_\eta \left(\frac{g_z}{g_\eta}\right)^{1/4} \left(\frac{\Lambda}{M_P}\right)^{13/2} \left(\frac{m_\eta}{m_{3/2}}\right)^{3/2}, \quad (3.14)$$

for the entropy ratio. Then for our nominal values of  $m_\eta \sim 10^{-5}M_P$  and  $m_{3/2} \sim 10^{-13}M_P$  the entropy ratio is approximately  $10^{11} \Lambda^{13/2}$ . Clearly for  $\Lambda \sim M_P$ , a huge amount of entropy is produced as in the original Polonyi scenario. However, as one can, the entropy increase is a sensitive function of the stabilization scale  $\Lambda$  and for  $\Lambda \lesssim 10^{-2}M_P$ , the entropy increase becomes tolerable.

For smaller  $\Lambda$ , even if  $Z$  decays after reheating, the energy density may never become dominated by  $z$  oscillations. In this case, the scale factor at the time of decay is such that  $\Gamma_z = t^{-1} = 2H_r$ . With the decay width given by eq. (2.24), the scale factor at  $Z$  decay,  $R_{dz}$  is found to be

$$\frac{R_{dz}}{R_\eta} \simeq 0.9 d_\eta^{-1/3} \Lambda^{5/2} m_\eta^{1/6} m_{3/2}^{-3/2} M_P^{-7/6}. \quad (3.15)$$

This assumes that the universe is dominated by radiation when  $Z$  decays, i.e.,  $\rho_r/\rho_z > 1$ . This is valid so long as the parameter  $\Lambda$  satisfies the constraint

$$\Lambda \lesssim 1.6 d_\eta^{-2/13} \left( \frac{m_{3/2}}{m_\eta} \right)^{3/13} M_P = 0.02 \tilde{d}_\eta^{-2/13} \left( \frac{m_{3/2}}{10^{-13} M_P} \frac{10^{-5} M_P}{m_\eta} \right)^{3/13} M_P. \quad (3.16)$$

If the limit in (3.16) is satisfied, the universe is never dominated by  $z$  oscillations, and  $\rho_z < \rho_r$  at the time of decay. In this case there will be no net entropy production. Therefore, for  $\Lambda \lesssim 10^{-2} M_P$ , all the cosmological problems associated with the evolution of the hidden sector are resolved. In particular, no significant amounts of entropy are generated, and any dilution effects of the products of baryogenesis and nucleosynthesis may be neglected.

Next, we return to the case that  $\Lambda$  is small enough so that  $z$  decay occurs before inflaton decay. The ratio  $R_{dz}/R_{d\eta}$  is smaller than one for

$$\Lambda \lesssim d_\eta^{-2/5} \left( \frac{m_{3/2}}{m_\eta} \right)^{3/5} M_P = 1.6 \times 10^{-5} \tilde{d}_\eta^{-2/5} \left( \frac{m_{3/2}}{10^{-13} M_P} \frac{10^{-5} M_P}{m_\eta} \right)^{3/5} M_P. \quad (3.17)$$

Thus, for smaller  $\Lambda$ , the decay of the Polonyi field occurs before reheating. In this scenario, the decay occurs when  $\Gamma_z = \frac{3}{2}H$ , with the Hubble parameter given by eq. (3.4). The scale factor is given by

$$\frac{R_{dz}}{R_\eta} = \frac{\pi^{2/3}}{3} \Lambda^{10/3} m_\eta^{2/3} m_{3/2}^{-2} M_P^{4/3}. \quad (3.18)$$

The universe is dominated by the oscillations of the inflaton field, since  $\rho_\eta/\rho_z = 16(\Lambda/M_P)^{-4} \gg 1$  at  $Z$  decay. The entropy release due to the modulus decay is clearly negligible in this case.

For completeness, we also consider the case where  $Z$  oscillations begin after inflaton decay. The scale factor  $R_z$  is found equating  $\frac{2}{3}m_z$  with the Hubble parameter  $H$  given by (3.10),

$$\frac{R_z}{R_\eta} = 0.5 d_\eta^{-1/3} \frac{\Lambda^{1/2} m_\eta^{1/6}}{m_{3/2}^{1/2} M_P^{1/6}}. \quad (3.19)$$

Then  $Z$  decay occurs when

$$\frac{R_{dz}}{R_\eta} = 0.9 d_\eta^{-1/3} \frac{\Lambda^{5/2} m_\eta^{1/6}}{m_{3/2}^{3/2} M_P^{7/6}}, \quad (3.20)$$

provided that  $Z$  decays before it dominates. Entropy production in this case is averted if  $\Lambda < 10^{-2} M_P (m_{3/2}/10^{-13} M_P)^{1/6}$ . Otherwise,  $R_{dz}$  is given by (3.13).

#### 4 Dark matter production and the gravitino problems

Having resolved the problem of entropy production, we turn to another of the serious issues facing moduli in cosmology, namely the overproduction of non-thermal relics. As we have

shown, the Polonyi modulus decays predominantly into a pair of gravitinos. This implies that the gravitino density produced by  $Z$  decay is  $n_{3/2} = 2n_z$ , where  $n_z$ ,  $n_{3/2}$  denote the number density of  $Z$  and the gravitino respectively. In addition, inflation may be a source of gravitinos through direct decay or through thermal processes during reheating. Gravitinos in turn will decay into an odd number of lightest supersymmetric particles (LSP), provided  $R$ -parity is a good symmetry. If the decay of  $Z$  into gravitinos is too efficient ( $\Lambda$  much smaller than one), or if the thermal reheat temperature after inflation is too high, gravitinos may be too copiously produced, and the resulting LSP abundance will be large enough to over-close the universe. We will assume that the LSP corresponds to a neutralino. In this case, direct production of LSPs by the decay of  $Z$  is negligible.

The measured cold dark matter density [39], assumed to be neutralinos, leads to a direct bound on the abundance of gravitinos. The closure fraction in neutralinos produced by gravitino decay, assuming  $n_\chi = n_{3/2}$ , can be written as

$$\Omega_\chi \simeq \frac{7m_\chi n_{3/2} n_\gamma}{s \rho_c} \simeq 2.75 \times 10^{10} h^{-2} \left( \frac{m_\chi}{100 \text{ GeV}} \right) \frac{n_{3/2}}{s}, \quad (4.1)$$

where  $s$  is the entropy density  $\simeq 7n_\gamma$  today and  $\rho_c$  is the closure density. Thus for  $\Omega_\chi h^2 \lesssim 0.12$ , we have an upper limit

$$\frac{n_{3/2}}{s} \lesssim 4.4 \times 10^{-12} \left( \frac{100 \text{ GeV}}{m_\chi} \right). \quad (4.2)$$

For sufficiently heavy gravitinos ( $m_{3/2} \gtrsim 10 \text{ TeV}$ ), this bound dominates over the limit from big bang nucleosynthesis (see e.g. [52, 53]).

Naively, the direct decay of an inflaton to a gravitino and inflatino could easily violate the bound (4.2). If one assumed that the density of gravitinos was equal or close to the number density of inflatons prior to their decay, the gravitino density would scale as  $n_\eta/s \sim (m_\eta/M_P)^{1/2}$ . However, direct decays into a gravitino and an inflatino may be kinematically forbidden [54] if  $|m_\eta - m_{\tilde{\eta}}| < m_{3/2}$  where  $m_{\tilde{\eta}}$  is the mass of the inflatino, or kinematically suppressed [55] if  $m_{3/2} \ll m_\eta \simeq m_{\tilde{\eta}}$ . In that case,  $n_{3/2}/s \sim (m_\eta/M_P)^{1/2} (m_{3/2}/m_\eta)$  and would safely satisfy the bound (4.2). Decays of the inflaton into a pairs of gravitinos can also be problematic [49]. This decay proceeds via the interaction  $G_\eta \eta \bar{\psi}_\mu \sigma^{\mu\nu} \psi_\nu$ . For a canonically normalized inflaton,  $G_\eta = \eta^\dagger$  and the decay rate depends on the vev of the inflaton after inflation. For models such as chaotic inflation [56] or Starobinsky models [57–59] where  $\langle \eta \rangle = 0$  after inflation, this decay mode will be effectively zero. We will therefore ignore the direct production of gravitinos from inflaton decay.

The thermal production of gravitinos during reheating is potentially more problematic as it is proportional to the reheat temperature (3.11). The gravitino-to-entropy ratio from thermal production is calculated to be [60–63]

$$\frac{n_{3/2}}{s} = 2.4 \times 10^{-12} \left( \frac{T_R}{10^{10} \text{ GeV}} \right) = 2.6 \times 10^{-11} d_\eta g_\eta^{-1/4} \left( \frac{m_\eta}{10^{-5} M_P} \right)^{3/2}, \quad (4.3)$$

for  $m_{1/2} \ll m_{3/2}$ . Combining eqs. (4.2) and (4.3) we have,

$$d_\eta g_\eta^{-1/4} \left( \frac{m_\eta}{10^{-5} M_P} \right)^{3/2} \left( \frac{m_\chi}{100 \text{ GeV}} \right) \lesssim 0.17, \quad (4.4)$$

which can clearly be satisfied.

Finally, we discuss the abundance of gravitinos computed by determining the number density of  $Z$  when it decays. As we will see, the limit on  $\Lambda$  from the non-thermal production of neutralinos is stronger than the limit from entropy production derived above. Therefore, in this section, we will assume that the bound (3.16) is satisfied and  $\Lambda$  is sufficiently small so that  $Z$  never dominates the energy density.

The number density of  $z$ 's is dependent on whether the decay occurs before or after reheating,

$$n_z = \frac{\rho_z}{m_z} \simeq \begin{cases} 0.033 d_\eta \Lambda^{-5/2} m_\eta^{3/2} m_{3/2}^{7/2} M_P^{1/2}, & R_{dz} > R_{d\eta} \\ 0.066 \Lambda^{-5} m_{3/2}^5 M_P^3, & R_{dz} < R_{d\eta} \end{cases} \quad (4.5)$$

In both cases, the resulting gravitino number density to entropy ratio is given by

$$\frac{n_{3/2}}{s} = 0.038 g_\eta^{-1/4} d_\eta \Lambda^5 \left( \frac{m_\eta^{3/2}}{m_{3/2} M_P^{11/2}} \right). \quad (4.6)$$

The corresponding neutralino yield is  $n_\chi/s \simeq n_{3/2}/s$ . Thus, the neutralino density parameter  $\Omega_\chi = m_\chi n_\chi / \rho_c$  is evaluated to be

$$\Omega_\chi h^2 \simeq 0.12 g_\eta^{-1/4} d_\eta \left( \frac{\Lambda}{8 \times 10^{-4} M_P} \right)^5 \left( \frac{m_\chi}{100 \text{ GeV}} \right) \left( \frac{m_\eta}{10^{-5} M_P} \right)^{3/2} \left( \frac{10^{-13} M_P}{m_{3/2}} \right), \quad (4.7)$$

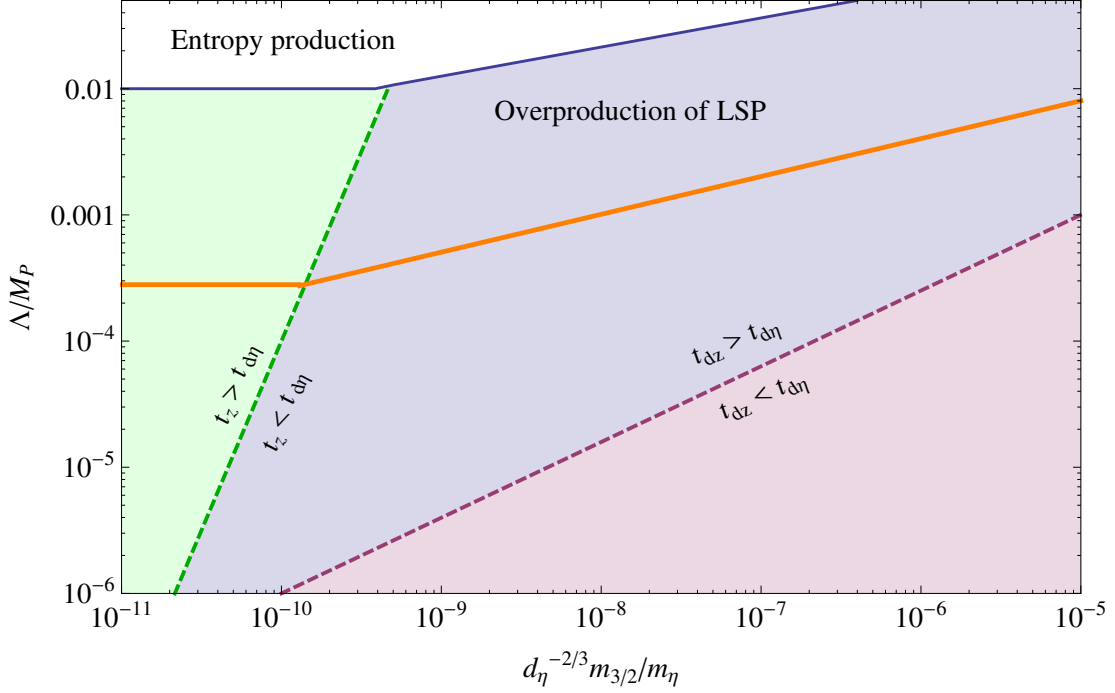
and the scale  $\Lambda$  which provides the necessary strong stabilization for the Polonyi modulus may be tuned to yield a density parameter consistent with the Planck normalization for the dark matter content of the universe [39]. The value of  $\Lambda$  for which the correct relic density is obtained corresponds to the reheating-before-decay scenario (see figure 2), and to the mass  $m_Z \sim 10^9 \text{ GeV}$ .

The validity of the expression (4.7) for the density parameter depends on the assumption that no significant amount of entropy is released at the decay of the gravitino. Since gravitinos are weakly interacting and non-thermally produced from the decay of the Polonyi field, they are effectively thermally decoupled until their decay. Gravitinos are relativistic at the time of production, since  $m_z \gg m_{3/2}$ , but they are slowed down by redshift, with momenta  $p \propto R^{-1}$  [64]. The dominant decays of the gravitino correspond to decays into a standard model particle and its supersymmetric partner, with rate<sup>1</sup> [65]

$$\Gamma_{3/2}(\psi_{3/2} \rightarrow \text{MSSM}) \simeq \frac{193}{384\pi} \frac{m_{3/2}^3}{M_P^2}. \quad (4.8)$$

With the scale factor at gravitino decay given by  $(R_{3/2}/R_{dz})^2 = \Gamma_z/\Gamma_{3/2} \simeq 10.3(\Lambda/M_P)^{-5} \gg 1$ , the gravitino will be non-relativistic at the time of decay. Approximating the energy density as  $\rho_{3/2} = \rho_z(m_{3/2}/2m_z)(R_{dz}/R)^3$ , with  $\rho_z$  the energy of the Polonyi field at its decay, a straightforward calculation shows that when the gravitino decays, the universe is dominated by the relativistic products of the inflaton,  $\rho_r/\rho_{3/2} > 1$ , if the mass scale  $\Lambda$  satisfies  $\Lambda \lesssim 8.2 \times 10^{-3} M_P$ . Therefore, the gravitino never dominates the universe, and no significant amount of entropy is released at its decay.

<sup>1</sup>If the sfermions are kinematically inaccessible to the gravitino, the decay rate is only marginally changed. The new rate is the same as that in eq. (4.8) except now take  $193 \rightarrow 144$ .



**Figure 2.** Allowed range for  $\Lambda$  as a function of  $d_\eta^{-2/3} m_{3/2}/m_\eta \simeq 1.3g^{-1/6} m_{3/2}/(T_R^2 M_P)^{1/3}$ . For the scenario in which  $Z$  oscillates before reheating, the upper limit corresponds to (3.16), while the boundary at which  $Z$  decays at reheating is given by (3.17). The orange curve corresponds to eq. (4.7), which can be rewritten as  $\Lambda/M_P = 0.25 (g/g_{\text{SM}})^{1/20} (\Omega_\chi h^2/0.1199)^{1/5} (m_\chi/100 \text{ GeV})^{-1/5} \cdot (m_{3/2}/10^{-13} M_P)^{-1/10} (d_\eta^{-2/3} m_{3/2}/m_\eta)^{3/10}$ .

Since the thermal production of gravitinos is independent of  $\Lambda$ , we can compare the thermal abundance with that produced by Polonyi decays. The ratio of the gravitino yield produced by modulus decay (4.6) to the thermally produced yield (4.3) is

$$\frac{(n_{3/2}/s)_{Z \text{ decay}}}{(n_{3/2}/s)_{\text{thermal}}} \simeq 4.58 \times 10^{14} \left( \frac{\Lambda}{M_P} \right)^5 \left( \frac{10^{-13} M_P}{m_{3/2}} \right). \quad (4.9)$$

Using this comparison and assuming that we satisfy the bound (4.4), we can obtain a bound on  $\Lambda$  which insures that the number of gravitinos produced by Polonyi decay is subdominant. This is the case if

$$\Lambda \lesssim 1.2 \times 10^{-3} M_P \left( \frac{m_{3/2}}{10^{-13} M_P} \right)^{1/5}. \quad (4.10)$$

Finally, if we include the effects of annihilations, the neutralino abundance produced by gravitino decay is determined from the Boltzmann equation

$$\frac{dn_\chi}{dt} + 3Hn_\chi = -\langle \sigma_{\text{ann}} v_{\text{rel}} \rangle n_\chi^2, \quad (4.11)$$

where  $\langle \sigma_{\text{ann}} v_{\text{rel}} \rangle$  denotes the thermal-averaged annihilation cross section. If the universe is dominated by the energy density of radiation,  $\rho_r > \rho_{\text{LSP}}$ , and since the entropy release from the gravitino decay is negligible, the relic abundance is found to be [50, 51, 66, 67]

$$\left( \frac{n_\chi}{s} \right)^{-1} \simeq \left( \frac{n_\chi}{s} \right)^{-1}_{3/2} + \left( \frac{H}{s \langle \sigma_{\text{ann}} v_{\text{rel}} \rangle} \right)^{-1}_{3/2}, \quad (4.12)$$

where the subindex indicates evaluation at gravitino decay. Therefore, the previous result (4.7) is only altered if the annihilation term is smaller than the gravitino yield (4.6). For typical annihilation rates for neutralino LSP,  $\langle\sigma_{\text{ann}}v_{\text{rel}}\rangle \sim 10^{-7} - 10^{-8} \text{ GeV}^{-2}$  [5, 67, 68], the ratio

$$\left(\frac{H/s\langle\sigma_{\text{ann}}v_{\text{rel}}\rangle}{n_{\chi}/s}\right)_{3/2} \sim \frac{10}{d_{\eta}} \left(\frac{\Lambda}{8 \times 10^{-4} M_P}\right)^{-5} \left(\frac{m_{\eta}}{10^{-5} M_P}\right)^{-3/2} \left(\frac{m_{3/2}}{10^{-13} M_P}\right)^{-1/2} \left(\frac{10^{-7} \text{ GeV}^{-2}}{\langle\sigma_{\text{ann}}v_{\text{rel}}\rangle}\right)$$

is larger than one, indicating that pair annihilation of neutralinos is not effective, and all the produced LSP's during gravitino decay survive.

## 5 Summary and conclusion

We have considered the cosmological consequences of a strongly stabilized, supersymmetry breaking hidden sector. The degree of stabilization is characterized by a mass scale,  $\Lambda$ , defined in the Kähler potential. We have shown that solutions to the cosmological problems inherent to light moduli in supergravity are possible for sufficiently small  $\Lambda$ . In this approach, the Polonyi sector is not responsible for providing the reheating temperature necessary for nucleosynthesis, since its energy density and the entropy released by its decay are subdominant with respect to that of the inflaton field. This restriction could easily be relaxed in some scenarios of the Affleck-Dine mechanism of baryogenesis where the late entropy release from the modulus decay is necessary to dilute a large baryon asymmetry. Nevertheless, a large baryon asymmetry is not a generic feature of the Affleck-Dine mechanism, and a negligible entropy release from modulus decay is in some cases necessary to obtain an asymmetry consistent with observations. This is true in particular when the flat direction responsible for the asymmetry is lifted by non-renormalizable quartic operators in the superpotential [13, 14, 36, 69].

Our results are neatly summarized in figure 2 which shows the various physical regimes discussed above for  $\Lambda$  as a function of the dimensionless combination  $d_{\eta}^{-2/3} m_{3/2}/m_{\eta}$ . At large  $\Lambda$ , there is an excessive amount of entropy produced as in the classic Polonyi scenario. At somewhat lower  $\Lambda$ , although the Polonyi field never comes to dominate the energy density of the universe, its decay leads to the over-production of the LSP. The figure also demarcates the values of  $\Lambda$  such that the Polonyi field decays before or after the inflaton. The figure does not show, however, the additional constraint (4.4) derived from thermally produced gravitino decay as this constraint is independent of the dynamics of the Polonyi sector. In the pale green region on the left of the figure,  $Z$  oscillations begin after inflaton decay. In that case, the entropy production and the LSP relic density are independent of  $m_{\eta}$ , and the curves plotted assume a fixed value of  $m_{3/2} = 10^{-13} M_P$ .

It must be emphasized that the introduction of the single stabilizing parameter  $\Lambda$  not only accounts for the solution of the entropy problems related to the Polonyi field, but it may also preclude the later onset of a gravitino and neutralino problem from moduli decay. Unless the LSP is copiously produced during inflaton decay or by scatterings in the primordial plasma, the suppression of all decay channels of the hidden sector relative to the gravitino channel imply that the bulk of the relic LSP density is generated from the decay of the gravitinos produced by the modulus decay. In this sense, the decay of the strongly stabilized Polonyi field can account for the present dark matter abundance. The constraint on  $\Lambda$  coming from the observed abundance,  $\Omega_{\chi} h^2 = 0.1199$ , lies well within the bound imposed by the resolution of the cosmological problems for the Polonyi field.

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